

More than mass proportional heating of heavy ions by supercritical collisionless shocks in the solar corona

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(Dated: June 17, 2009)

Abstract

We propose a new model for explaining the observations of more than mass proportional heating of heavy ions in the polar solar corona. We point out that a large number of small scale intermittent shock waves can be present in the solar corona. The energization mechanism is, essentially, the ion reflection off supercritical quasi-perpendicular collisionless shocks in the corona and the subsequent acceleration by the motional electric field $\mathbf{E} = -(1/c)\mathbf{V} \times \mathbf{B}$. The acceleration due to \mathbf{E} is perpendicular to the magnetic field, in agreement with observations, and is more than mass proportional with respect to protons, because the heavy ion orbit is mostly upstream of the quasi-perpendicular shock foot. The observed temperature ratios between O^{5+} ions and protons in the polar corona, and between α particles and protons in the solar wind are easily recovered.

The heating of the solar corona to temperatures of the order of 10^6 K and more is one of the outstanding problems of solar physics. Beside the high temperatures, Soho/UVCS observations have shown that heavy ions in polar corona, like O^{5+} and Mg^{9+} , are heated more than protons, and that heavy ion heating is more than mass proportional; further, the perpendicular temperatures T_{\perp} are much larger than parallel temperatures T_{\parallel} [1, 2, 3]. As a consequence of magnetic mirroring in the diverging magnetic field of coronal holes, heavy ions are observed to be faster than protons in the solar wind [2, 4]. In addition, the collisional coupling with protons up to $1.32 R_{\odot}$ indicates that the Mg^{9+} heating has to be faster than minutes [5]. These observations give stringent constraints on the coronal heating mechanism. Ion cyclotron heating has been considered since long (e.g., [3, 6, 7, 8]), but some details are not yet fully understood. The comprehension of coronal heating is of general physical interest, as the sun serves both as a huge plasma laboratory and as a model for a large class of stars.

Shock waves are considered to be common in the chromosphere/transition region and in the corona (e.g., [9, 10, 11, 12, 13]). For instance, photospheric convection leads to the emergence of small magnetic loops, which lead to magnetic reconnection with the network magnetic field; small scale plasma jets are formed in the reconnection regions, and fast shocks can form when jets encounter the ambient plasma [10, 11, 14]. Indeed, recent X-ray Hinode and UV Stereo observations have shown that many more plasma jets are present in the polar corona than previously thought [15, 16, 17]. Therefore, a large number of small scale, intermittent shocks can form in the reconnection regions and propagate toward the high altitude corona. Recent numerical simulations show that bursty, time dependent reconnection in solar flares can eject many plasmoids and create oblique shocks [18]. In the high corona, magnetic reconnection happens when current sheets form because of the evolving coronal structures [19], while large scale shocks propagate in the corona because solar flares and of the emergence of coronal mass ejections [20, 21]. Such shocks are detected as type II radio bursts [22]. In the case of solar flares, the associated reconnection outflow termination shocks can be so strong as to accelerate electrons to 100 keV energies in a fraction of a second [14, 23], and in some cases both the upper and the lower termination shocks are identified in radio observations [24].

In the low β , nearly collisionless corona, a shock wave is formed when a superAlfvénic plasma flow having velocity $V_1 > V_A$ collides with the ambient coronal plasma. Here, the

plasma β is given by $\beta = 8\pi p/B^2$, where p is the total plasma pressure, B is the magnetic field magnitude, V_1 is the plasma velocity upstream of the shock, and $V_A = B/\sqrt{4\pi\rho}$ is the Alfvén velocity, with ρ the mass density. The Alfvénic Mach number is defined as $M_A = V_1/V_A$. We notice that although the typical Alfvén velocity in the corona, of the order of 1000 km/s, is larger than the observed jet velocities of 200–800 km/s [15, 25], the Alfvén velocity in the reconnection region can be much lower, since B is weaker there. Indeed, the reconnection regions are characterized by current sheets, magnetic field reversals, and magnetic quasi-neutral lines. For instance, Tsuneta and Naito [14] argue that an oblique fast shock is naturally formed below the reconnection site in the corona, see their Figure 1. The plasma velocity in the reconnection outflow region between the slow shocks is of the order of V_A in the inflow region, that is much larger than V_A in the outflow region, and this leads to the formation of shocks (e.g., [9, 14, 20]).

Previously, shock heating of coronal heavy ions was considered by Lee and Wu [11], but mostly in connection with subcritical shocks. Here, we propose that the more than mass proportional heating of heavy ions in polar coronal holes is due to ion reflection at supercritical quasi-perpendicular shocks and to the ion acceleration by the $\mathbf{V} \times \mathbf{B}$ electric field in the shock frame. In this connection, we notice that more than mass proportional heating of α particles and O^{6+} has been observed in the solar wind by the Ulysses spacecraft, between 2.7 and 5.1 AU, downstream of interplanetary shocks, most of which were supercritical [26] (see also Ref. [27]).

It is well known both from laboratory [28, 29] and from spacecraft experiments (e.g., [30, 31, 32]) that above a critical Mach number $M_A^* \simeq 2.7$ for perpendicular collisionless shocks (less than 2.7 if the shock is quasi-perpendicular), a fraction of ions, which grows with the Alfvénic Mach number [29, 33], is reflected off the shock, leading to the so-called supercritical shocks. When the angle θ_{Bn} between the shock normal (pointing in the upstream direction) and the upstream magnetic field is larger than about 45° , the reflected ions reenter the shock after gyrating in the upstream magnetic field. Such shocks are termed quasi-perpendicular. Conversely, for $\theta_{Bn} < 45^\circ$, the reflected ions propagate upstream, forming the ion foreshock which characterizes the quasi-parallel shocks. The critical Mach number M_A^* can decrease below 1.5 for oblique shocks in a warm plasma [34], so that ion reflection is a relatively common process. Ion reflection can be considered to be the main dissipation mechanism by which collisionless shocks convert the flow directed energy into heat, while the electrons are

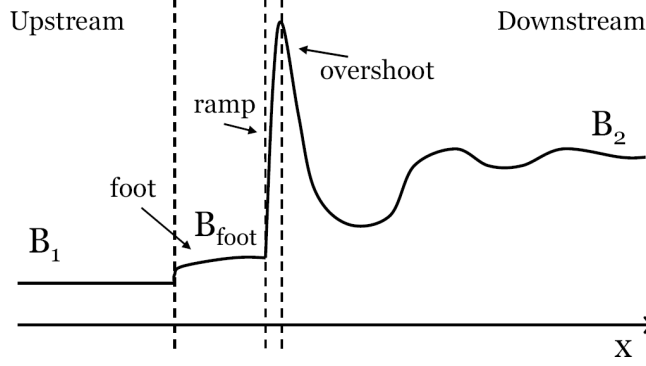


FIG. 1: Schematic of the magnetic field profile of a supercritical quasi-perpendicular collisionless shock. The main features like the magnetic foot, the ramp, and the magnetic overshoot are indicated.

heated much less (typically, one tenth of proton heating) [31].

For the solar corona, we consider a quasi-perpendicular supercritical collisionless shock, and we assume a simple one dimensional shock structure. The upstream quantities are indicated by the subscript 1, and the downstream quantities by the subscript 2. We adopt the Normal Incidence Frame (NIF) of reference, in which the shock is at rest, the upstream plasma velocity is along the x axis and perpendicular to the shock surface, $\mathbf{V}_1 = (V_{x1}, 0, 0)$, the upstream magnetic field lays in the xz plane, $\mathbf{B}_1 = (B_{x1}, 0, B_{z1})$, so that the motional electric field $\mathbf{E} = -\mathbf{V} \times \mathbf{B}/c$ is in the y direction, $E_y = V_{x1}B_{z1}/c$. We further assume that $B_{z1} \gg B_{x1}$ ($\theta_{Bn} \simeq 90^\circ$), in order to simplify the discussion. An order-of-magnitude estimate of the energy gained by ions after reflection at the shock can be obtained by approximating the reflected ion trajectory with a circle of radius r_L , with r_L the ion Larmor radius, in the upstream magnetic field. Assuming specular reflection [30, 31, 35], on average the ion velocity at the reflection point is perpendicular to the shock and along the x axis. Keeping in mind that ion reflection gives rise to a non adiabatic displacement in the y direction, the work W done by the electric field is

$$W = q_i E_y \Delta y, \quad (1)$$

where $\Delta y \sim 2r_L$. For specularly reflected ions, the Larmor radius has to be evaluated with the upstream flow speed (neglecting the thermal velocity of the incoming ion distribution), i.e., $v_\perp \simeq |V_{x1}|$, so that

$$W = q_i E_y \times 2r_L = 2q_i \frac{V_{x1}B_{z1}}{c} \frac{m_i V_{x1}c}{q_i B_{z1}}, \quad (2)$$

which yields $W = 2m_i V_{x1}^2$. This estimate shows that the energy gain is mass proportional.

A more detailed calculation yields a more precise result, and shows that heavy ion heating is more than mass proportional. In order to do this, we remind that a distinctive feature of quasi-perpendicular collisionless shocks is the formation of a “foot” in the magnetic field profile in front of the main magnetic ramp, the latter culminating in the magnetic overshoot, beyond which the downstream values are gradually attained [29, 32, 36], see Figure 1. The foot is due to the population of reflected and gyrating protons, which causes an increase in the plasma density, and, because of the magnetized electrons, in the magnetic field strength [29]. Even if the solar corona composition encompasses several ion species, the foot extent in the upstream direction is determined by the proton gyroradius, since protons are the major species. We define $B_{\text{foot}} = (1 + b)B_{z1}$, with b depending on the ion reflection rate; b can be estimated to be of the order 0.5–1 for typical shocks in the heliosphere [31, 32, 36]. Direct observations in space show that very strong fluctuation levels are found in association with collisionless shocks. However, in what follows we will neglect fluctuations and we will consider only the average quantities, in order to set the stage. Taking into account the fact that the orbit in crossed electric and magnetic fields is a trochoid, we start from the equations of the particle trajectory. We assume the magnetic field to be along the z axis, and set the origin of coordinates at the point of ion reflection, with $t = 0$ (e.g., Ref. [37]):

$$x(t) = -\frac{v_{\perp}}{\Omega} \sin(\Omega t) + \frac{cE_y}{B} t \quad (3)$$

$$y(t) = \frac{v_{\perp}}{\Omega} [1 - \cos(\Omega t)] \quad (4)$$

where $\Omega = q_i B / m_i c$, and B the local magnetic field. The corresponding particle velocity is

$$v_x(t) = -v_{\perp} \cos(\Omega t) + \frac{cE_y}{B} \quad (5)$$

$$v_y(t) = v_{\perp} \sin(\Omega t). \quad (6)$$

Specular ion reflection implies that at $t = 0$ the ion velocity v_x is opposite to the incoming plasma velocity, $v_x(t = 0) = -V_{x1}$, whence

$$v_{\perp} = V_{x1} + \frac{cE_y}{B_{\text{foot}}} = V_{x1} + \frac{V_{x1} B_{z1}}{B_{\text{foot}}} = V_{x1} \frac{2 + b}{1 + b}. \quad (7)$$

The reflected ions meet again the shock surface, at $x = 0$, at a later time $t_1 > 0$, corresponding to

$$\frac{v_{\perp}}{\Omega} \sin(\Omega t_1) = \frac{cE_y}{B} t_1. \quad (8)$$

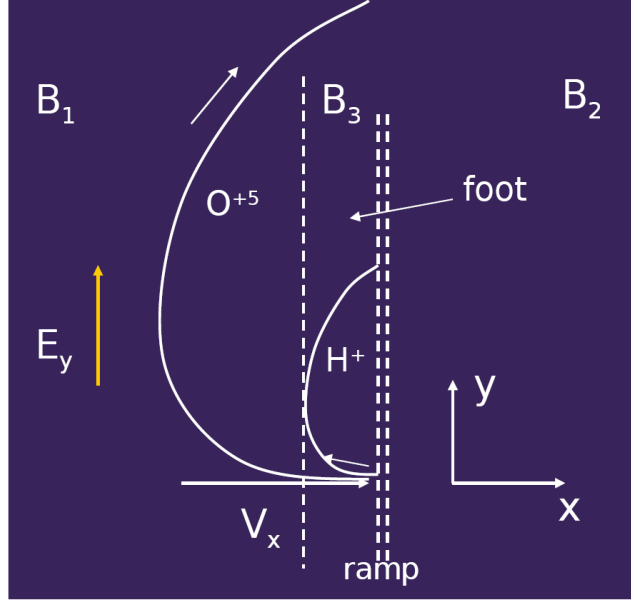


FIG. 2: (Color online) Projection in the xy plane of the trajectories of hydrogen and oxygen ions reflected at the shock ramp. The motional electric field is also indicated. As in Figure 1, the vertical dashed lines separate the main magnetic field regions, such as the upstream region, the foot, the ramp, and the downstream region.

Upon inserting the values of v_{\perp} and of E_y in the above equation we obtain

$$\sin(\Omega t_1) = \frac{\Omega t_1}{2 + b}, \quad (9)$$

whose numerical inversion yields $\Omega t_1 = 2.27885$ for $b = 1$, $\Omega t_1 = 2.1253$ for $b = 0.5$, and $\Omega t_1 = 1.8955$ for $b = 0$ (see below). At this time the particle will have moved in the y direction by an amount given by

$$\Delta y(t_1) = \frac{v_{\perp}}{\Omega} [1 - \cos(\Omega t_1)] = \frac{m_i V_{x1} c}{q_i B_{z1}} [1 - \cos(\Omega t_1)] \frac{2 + b}{(1 + b)^2}. \quad (10)$$

This displacement in the y direction determines the energy gained by reflected ions during the gyromotion in the field E_y :

$$W = q_i E_y \Delta y = m_i V_{x1}^2 [1 - \cos(\Omega t_1)] \frac{2 + b}{(1 + b)^2} = \frac{2(2 + b)}{(1 + b)^2} [1 - \cos(\Omega t_1)] \frac{1}{2} m_i V_{x1}^2 \quad (11)$$

Taking into account the fact that protons move in the foot magnetic field B_{foot} , we can assume that $b \simeq 0.5$ – 1 . In such a case, $1 - \cos(\Omega t_1) = 1.65035$ for $b = 1$ (or 1.52652 for $b = 0.5$), so that

$$W_p \simeq \frac{3}{2} \times 1.65035 \times \left(\frac{1}{2} m_p V_{x1}^2 \right). \quad (12)$$

On the other hand, for heavy ions like O^{5+} most of the trajectory is upstream of the foot, see Fig. 2, in the unperturbed plasma where $B \simeq B_{z1}$. Then we can set $b = 0$ with good approximation, and obtain $1 - \cos(\Omega t_1) = 1.319$, so that

$$W_{\text{heavy}} \simeq 4 \times 1.319 \times \left(\frac{1}{2} m_i V_{x1}^2\right). \quad (13)$$

Here we can see that, with respect to protons, heating is more than mass proportional. For $b = 1$, the ratio of the heavy ion energy gain over the proton energy gain is $W_{\text{heavy}}/W_p \simeq 2.13 \times m_i/m_p$, while for $b = 0.5$ we find $W_{\text{heavy}}/W_p \simeq 1.55 \times m_i/m_p$. Varying the value of b between 0.5 and 1 yields an O^{5+} temperature about 25–34 times larger than the proton temperature ($m_O \simeq 16m_p$), in good agreement with Soho/UVCS observations which give $T_{O^{5+}}/T_p = 27\text{--}37$ [5]. Also, assuming a typical value of $b = 0.5$, we can easily recover the temperature ratios observed in the solar wind for helium ($m_\alpha \simeq 4m_p$), where $T_\alpha/T_p \simeq 6$ is observed in those solar wind periods which are found to be the least collisional [38], thus reflecting more closely the coronal conditions. At the same time, for $b = 1$, $T_\alpha/T_p \simeq 8$ is obtained, a value which is also observed and reported in Figure 1 of Ref. [38].

On the other hand, heating is essentially perpendicular, since it is due to the motional electric field E_y which is perpendicular to the magnetic field by definition. This allows to understand the observed strong temperature anisotropy with $T_\perp \gg T_\parallel$. In addition, a single shock encounter is required to accelerate the ions, and the acceleration time is on the scale of the ion gyroperiod, so that the heating mechanism is very fast, as required by the observations in Ref. [5].

Typically, the reflection rate for protons is found from numerical simulations to be 20–30%. The energy gained by reflected ions is distributed to the transmitted ions by wave particle interactions [31], so that we can assume for the bulk of protons a heating rate about 1/4 of W_p . Let us define the heating efficiency, η , for protons, as the ratio of the energy gain of both reflected and transmitted ions over the upstream thermal energy $\frac{1}{2}m_p v_{th1}^2$:

$$\eta = \frac{W_p/4}{\frac{1}{2}m_p v_{th1}^2} \simeq \frac{2.48(\frac{1}{2}m_p V_{x1}^2)}{4 \times \frac{1}{2}m_p v_{th1}^2} = 0.62 M_s^2 \quad (14)$$

where $M_s = V_{x1}/v_{th1}$ is the sonic Mach number. In the low β corona, the thermal speed is much less than the Alfvén speed, so that $M_s \gg M_A$. However, in the reconnection outflow region the magnetic field is weaker than in the ambient corona, so that for a first estimate we assume that the sonic Mach number is of the same order of the Alfvén Mach number. For

instance, assuming that $M_s = 7$, we can see that the efficiency for protons equals $\eta \simeq 30$. In other words, the crossing of two or three shocks might bring the chromospheric plasma from temperatures of the order of 10^4 K to coronal temperatures of the order of 10^6 K. As a general trend, we can say that the stronger the magnetic field in the reconnection inflow region, the larger the plasma velocity and the Mach numbers in the reconnection outflow region, and the larger the heating efficiency. On the other hand, several shock crossings may be required for the high altitude corona to reach the observed temperatures, and we notice that large scale shocks associated with coronal mass ejections and type II radio bursts can propagate all the way into polar corona.

Clearly, the present model has to be further developed, since a wide range of different Mach numbers, plasma β , shock normal angles θ_{Bn} , fluctuation levels, and heating ratios can be envisaged in the corona. In addition, collisional coupling with protons can decrease the obtained temperature ratios, as may be the case of Mg^{9+} . Further, multiple reflected ions at a single shock can also be envisaged [43], a phenomenon which would increase the heating efficiency.

Our model leads to the prediction that a fraction of heavy ions comparable to the proton fraction is also reflected at quasi-perpendicular shocks. In collisionless shocks, an electrostatic potential barrier $\Delta\phi \simeq m_p(V_{x1}^2 - V_2^2)/2e$ arises, which slows down the incoming ions [29, 39, 44]. In simple discussions of ion reflection, ions are expected to undergo specular reflection if their kinetic energy in the shock frame is less than the potential energy barrier $q_i\Delta\phi$: however, for heavy ions this is usually found only for a tiny fraction of the upstream velocity distribution. Nevertheless, experimental evidence of α particle reflection at Earth's quasi-perpendicular bow shock has been reported by in Ref. [40], while evidence of α particle specular reflection off the quasi-parallel bow shock was reported in Ref. [41]. On the other hand, laboratory experiments show that increasing M_A , the number of reflected ions increases while the potential jump *decreases* [29], contrary to expectations if the potential jump would be the only cause of ion reflection. This shows that ions are not simply reflected by the average potential jump across the shock, and that also the magnetic foot and the fluctuating electric and magnetic overshoots play a role for ion reflection [30, 35, 39, 42]. Indeed, cross shock electric fields measured by the Polar spacecraft at Earth's bow shock show that the potential $\Delta\phi$ is strongly spiky and fluctuating [44]. Further, quasi-perpendicular shocks also exhibit cyclic reformation, which implies time-depending electric and magnetic

overshoots and reflection rate [33, 45]. Indeed, the Cluster spacecraft have recently shown that ion reflection is highly unsteady [45], so that the strong variations in the shock structure can also induce heavy ion reflection.

In conclusion, we have considered the heavy ion energization due to the ion reflection off quasi-perpendicular shocks. Fast, supercritical shocks are assumed to form because of reconnection of small scale magnetic loops at the base of coronal holes, like those associated with polar coronal jets, and because of the merging of magnetic structures in the higher corona. The energy stored in the coronal magnetic field is transformed to bulk kinetic energy by reconnection, and into heat and heavy ion heating by the quasi-perpendicular shocks which form in the reconnection outflow region. Our model can explain both coronal heating and the more than mass proportional heavy ion heating observed by Soho/UVCS. In addition, this heating mechanism is strictly perpendicular to the magnetic field and it is very fast (a single shock encounter is needed); most heating goes into the ions, with electrons undergoing an energy gain which is about an order of magnitude smaller than that of protons [31, 32]. Further, the strongly anisotropic heating with $T_{\perp} \gg T_{\parallel}$ can give rise to efficient ion cyclotron emission; this phenomenon is actually observed downstream of the Earth's bow shock. Indeed, recent Stereo, Helios, and Venus Express data show that ion cyclotron waves are probably present in the corona [46]. Therefore, quasi-perpendicular collisionless shocks can be the source of ion cyclotron waves in the corona, too. These waves later on can heat locally the solar wind by cyclotron resonance dissipation, as suggested by a number of observations.

This research was partially supported by the Italian INAF and the Italian Space Agency, contract ASI n. I/015/07/0 “Esplorazione del Sistema Solare”.

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